

# A Randall-Sundrum scenario with bulk dilaton and torsion

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## Abstract

We consider a string-inspired torsion-dilaton-gravity action in a Randall-Sundrum brane world scenario and show that, in an effective four-dimensional theory on the visible brane, the rank-2 antisymmetric Kalb-Ramond field (source of torsion) is exponentially suppressed. The result is similar to our earlier result in [13], where no dilaton was present in the bulk. This offers an explanation of the apparent invisibility of torsion in our space-time. However, in this case the trilinear couplings  $\sim \text{TeV}^{-1}$  between dilaton and torsion may lead to new signals in TeV scale experiments, bearing the stamp of extra warped dimensions.

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# 1 Introduction

Braneworld models with extra spatial dimensions have gained considerable interest over the last decade, largely due to their role in solving the naturalness/hierarchy problem in the context of the standard model of elementary particles. Extra spacelike dimensions with warped geometry elegantly bridges the gap of sixteen orders of magnitude between electroweak and the Planck scale. The key feature of the Randall-Sundrum type of theory[1] is the ‘ $-$ ’ exponential warp factor attached to the Minkowski part of the metric, leading to a suppression factor for all masses and gravitational couplings of known fields residing on the visible brane located at one of the orbifold fixed points of the extra compact spacelike dimension.

The essential input of the RS[1] type of theory is that gravity propagates in a 5-dimensional anti-de Sitter bulk spacetime, while the standard model (SM) fields are assumed to lie on a 3-brane. String theory provides a natural explanation for such description as the SM fields arise as excitations of an open string whose ends are fixed on the brane. The graviton, on the other hand, is a closed string excitation and hence can propagate in the bulk space-time[3]. Consequently, when one takes its projections on the visible brane, the massless graviton mode has a coupling  $\sim 1/M_P$  with all matter, while the massive modes have enhanced coupling through the warp factor. The massless graviton mode therefore accounts for the presence of gravity in our universe, while the massive modes raise hopes for possible new signals of extra warped dimension in accelerator experiments [4]. However, in a string-based model, one should also consider other closed string massless excitations which appear as various higher rank tensor fields in the low energy effective theory[3]. They are expected to propagate in the bulk just as gravity. Two such massless closed string excitations are scalar dilaton and second rank antisymmetric Kalb Ramond (KR) field [15]. The question we ask here is: can these fields have any observable effects? If not, why are the effects of their massless modes less perceptible on our brane than the force of gravitation? Implications of other type of bulk fields such as scalars[5], fermions [6] and gauge fields[7] have also been considered in this scenario. While some such scenarios are testable in accelerator experiments [8] or in the neutrino sector [9], by and large they do not cause any contradiction with our observations so far.

The antisymmetric rank-2 tensor bulk field namely the Kalb-Ramond[10] field has similar coupling to matter as gravity. Using a generalised form of the Einstein-Cartan action, it has been shown earlier that the third rank field strength tensor of KR field can be identified with torsion [11]. Experimental limits impose severe constraints on such fields so that it is vanishingly small[12]. This apparent invisibility of torsion has been clarified in one of our earlier work [13]. There it has been shown that the zero mode of the antisymmetric tensor

field gets an additional exponential suppression compared to the graviton on the visible brane. This could well be an explanation of why we see the effect of curvature but not of torsion in the evolution of the universe. But as mentioned earlier, string theory allow dilaton as well to exist in the bulk and hence it is essential to consider the dilaton-torsion coupling in the bulk, instead of taking a free torsion field. In this light, the earlier conclusion regarding the suppression of the torsion field on our brane must be re-examined when it is coupled to the dilaton in the bulk. We address this question in the current study in context of RS braneworld scenario. We also explore other possible implications of bulk dilaton field on the Physics on our brane. It may be mentioned that the issue of braneworld stability in a similar dilaton-torsion coupled scenario has been investigated by others[16].

We organise our work as follows. In the next section, we review the RS theory and then illustrate our model in the background of RS spacetime, where torsion coexist with gravity and dilaton in the bulk and has a non-minimal coupling with dilaton as one obtains in string inspired models. Section 3 presents the solution of the bulk dilaton on the visible brane. In section 4, we describe in detail how the different KK modes of the KR field are coupled to the dilaton fields on the brane, and how these couplings affect our previous explanation of the invisibility of torsion on our SM brane. From this, we conclude that even in presence of dilaton coupling the zero mode of the KR field is suppressed with an additional exponential suppression. Thus this work reinforces our claim that RS scenario can provide an explanation of why the effect of torsion is much weaker than that of curvature on the brane. However, because of the presence of the dilaton, we obtain trilinear couplings between dilaton-KR field which may have their imprints in Tev-scale Physics.

## 2 Dilaton-Torsion in Randall Sundrum framework

In the Randall-Sundrum model, we have a 5-dimensional anti de-sitter spacetime with the extra spatial dimension orbifolded as  $S_1/Z_2$ . Two 4 dimensional branes known as the visible (TeV) brane and hidden (Planck) brane are placed at the two orbifold fixed points  $y = 0$  and  $\pi$ . The 5-dimensional spacetime is characterised by the metric

$$ds^2 = e^{-2\sigma} \eta_{\mu\nu} dx^\mu dx^\nu + r_c^2 dy^2 \quad (1)$$

where  $\eta_{\mu\nu}$  is the usual 4D Minkowski metric with the sign convention  $(-, +, +, +)$ , and  $\sigma = kr_c|y|$ .  $r_c$  is the compactification radius for the fifth dimension, and  $k$  is on the order of the higher dimensional Planck mass  $M$ . (Greek indices are used for brane coordinates and Latin indices are used for the full 5D coordinates). The standard model fields reside at  $y = \pi$  while gravity peaks at  $y = 0$ . The dimensionful parameters defined above are related

to the 4-dimensional Planck scale  $M_P$  through the relation

$$M_P^2 = \frac{M^3}{k} [1 - e^{-2kr_c\pi}] \quad (2)$$

Clearly,  $M_P$ ,  $M$  and  $k$  are all of the same order of magnitude. As a consequence of warped geometry, all mass scales get exponentially lowered from the Planck scale to the TeV scale because of the warp factor  $e^{kr_c\pi}$ . Thus the hierarchy between the Planck and TeV scales can be explained without the need of any fine-tuning.

In the scenario adopted by us the source of torsion is taken to be the rank-2 anti symmetric KR field  $B_{MN}$  [11]. It has been shown in [11], that in the low energy effective string action the torsion field  $T_{MNL}$ , which is an auxiliary field, could be identified with the rank-3 anti symmetric field  $H_{MNL}$  by using the equation of motion  $T_{MNL} = H_{MNL}$ . Torsion can be identified with  $H_{MNL}$  which is the field strength tensor of the KR field  $B_{MN}$ .

In the context of string theory, apart from gravity, scalar dilaton and the second rank anti-symmetric tensor KR field both co-exist in the bulk. The five dimensional action for the gravity-dilaton-torsion sector is given by

$$S_{tot} = \int d^5x \sqrt{-G} [2M^3 R - e^{\phi/M^{3/2}} 2H_{MNL} H^{MNL} + \frac{1}{2} \partial_M \phi \partial^M \phi - m^2 \phi^2] \quad (3)$$

where

$$S_{grav} = \int d^5x \sqrt{-G} 2M^3 R \quad (4)$$

$$S_{KR} = \int d^5x \sqrt{-G} [-e^{\phi/M^{3/2}} 2H_{MNL} H^{MNL}] \quad (5)$$

$$S_{dil} = \int d^5x \sqrt{-G} [\frac{1}{2} \partial_M \phi \partial^M \phi - m^2 \phi^2] \quad (6)$$

From now onwards, we focus on the the part of the action consisting of  $S_{KR}$  and  $S_{dil}$  only. Our aim is to find out the influence of the dilaton field on the KR Lagrangian, and to see how it affects the existence of different torsion modes on the visible brane. Moreover, we take into account the massless dilaton field only, as the massive field on the bulk will anyway decouple. The relevant part of the action is,

$$S = \int d^5x \sqrt{-G} [-e^{\phi/M^{3/2}} 2H_{MNL} H^{MNL} + \frac{1}{2} \partial_M \phi \partial^M \phi] \quad (7)$$

Using the explicit form of the background metric in the Randall-Sundrum scenario and also the KR gauge fixing condition  $B_{4\mu} = 0$ , the above action can be written as

$$\begin{aligned} S = \int d^4x dy r_c [ & - 2e^{\phi/M^{3/2}} \{ e^{2\sigma} \eta^{\mu\alpha} \eta^{\nu\beta} \eta^{\lambda\gamma} \partial_{[\mu} B_{\nu\lambda]} \partial_{[\alpha} B_{\beta\gamma]} - \frac{3}{r_c^2} \eta^{\mu\alpha} \eta^{\nu\beta} B_{\mu\nu} \partial_y^2 B_{\alpha\beta} \} \\ & + \frac{1}{2} \{ e^{-2\sigma} \eta^{\mu\nu} \partial_\mu \phi \partial_\nu \phi - \frac{1}{r_c^2} \phi \partial_y [e^{-4\sigma} \partial_y \phi] \} ] \end{aligned} \quad (8)$$

Next, consider Kaluza-Klein decomposition of the following form for both the KR and the dilatonic field

$$B_{\mu\nu} = \sum_{n=0}^{\infty} B_{\mu\nu}^n(x) \frac{\chi^n(y)}{\sqrt{r_c}} \quad (9)$$

$$\phi = \sum_{n=0}^{\infty} \phi^n(x) \frac{\psi^n(y)}{\sqrt{r_c}} \quad (10)$$

In terms of the projections  $(B_{\mu\nu}^n(x), \phi^n(x))$  on the visible brane our effective action looks like

$$\begin{aligned} S = \int d^4x & \left[ \sum_n \sum_m \int dy e^{2\sigma} \chi^n \chi^m \left\{ -2e^{\frac{1}{M^{3/2}}} \sum_i \frac{\phi^i \psi^i}{\sqrt{r_c}} \eta^{\mu\alpha} \eta^{\nu\beta} \eta^{\lambda\gamma} \partial_{[\mu} B_{\nu\lambda]}^n \partial_{[\alpha} B_{\beta\gamma]}^m \right. \right. \\ & \quad \left. \left. - 3 \left[ \frac{1}{r_c^2} \frac{e^{-2\sigma}}{\chi^m} \partial_y^2 \chi^m \right] \eta^{\mu\alpha} \eta^{\nu\beta} B_{\mu\nu}^n B_{\alpha\beta}^m \right\} \right] \\ & + \frac{1}{2} \left[ \sum_n \sum_m \int dy e^{-2\sigma} \psi^n \psi^m \left\{ \eta^{\mu\nu} \partial_\mu \phi^n \partial_\nu \phi^m - \left[ \frac{1}{r_c^2} \frac{e^{2\sigma}}{\psi^m} \partial_y (e^{-4\sigma} \partial_y \psi^m) \right] \phi^n \phi^m \right\} \right] \end{aligned} \quad (11)$$

This essentially represents the four dimensional action where the term in the first square bracket is the lagrangian corresponding to the KR field coupled to dilaton and the second square bracketed term is the lagrangian for a dilaton field. It is interesting to note that the second term which is basically  $S_{dil}$  has no coupling with the KR field and hence corresponds to a free dilaton action. We first proceed with this part to find out how the massless bulk dilaton field that appears on the standard model brane.

### 3 Bulk Dilaton Field

It is straightforward to check from above that the dilaton action assumes a standard canonical form on the standard model brane i.e.,

$$S_{dil} = \int d^4x \sum_{n=0}^{\infty} \frac{1}{2} (\eta^{\mu\nu} \partial_\mu \phi^n \partial_\nu \phi^n + m_n^{d^2} \phi^{n^2}) \quad (12)$$

provided  $m_n^d$ , the mass of the nth mode of the dilaton field, is defined by

$$-\frac{1}{r_c^2} \frac{e^{2\sigma}}{\psi^m} \partial_y (e^{-4\sigma} \partial_y \psi^m) = m_n^{d^2} \quad (13)$$

and  $\psi^m(y)$  satisfy the orthonormality relation

$$\int_{-\pi}^{\pi} dy e^{-2\sigma} \psi^n(y) \psi^m(y) = \delta_{nm}. \quad (14)$$

In terms of  $z_n = \frac{m_n^d}{k} e^\sigma$  and  $f^n = \psi^n e^{-2\sigma}$  equation (13) can be recast in the form

$$z_n^2 \frac{d^2 f^n}{dz_n^2} + z_n \frac{df^n}{dz_n} + [z_n^2 - 4] f^n = 0 \quad (15)$$

The above equation admits a solution of Bessel function of order 2. Consequently,  $\psi^n$  has a solution of the form

$$\psi^n = \frac{e^{2\sigma}}{N_n} [J_2(z_n) + \alpha_n Y_2(z_n)] \quad (16)$$

where  $J_2(z_n)$  and  $Y_2(z_n)$  are Bessel and Neumann functions of order 2,  $\alpha_n$  and  $N_n$  are two constants to be determined from the boundary conditions at the orbifold fixed points. The continuity condition for the derivative of  $\psi^n$  at  $y = 0$ , dictated by the self-adjointness of the left hand side of equation(13), yields

$$\alpha_n = -\frac{J_1(\frac{m_n^d}{k})}{Y_1(\frac{m_n^d}{k})} \quad (17)$$

Using the fact that  $e^{kr_c\pi} \gg 1$  and the mass values  $m_n^d$  on the brane is expected to be on the order of TeV scale ( $\ll k$ ),

$$\alpha_n = -\frac{\pi}{4} \left(\frac{m_n^d}{k}\right)^2 \ll 1 \quad (18)$$

The boundary condition at  $y = \pi$  gives

$$J_1(x_n) = 0 \quad (19)$$

where  $x_n = z_n(\pi) = \frac{m_n^d}{k} e^{kr_c\pi}$ .

The roots of the above equation determine the masses of the higher excitation modes. As  $x_n$  is of order unity, the massive modes lie in the TeV scale. The point to note here is that the masses of the higher excitation modes of the dilaton field are same as the massive graviton modes.

The orthonormality condition determines the constant  $N_n$ :

$$N_n = \frac{e^{kr_c\pi}}{\sqrt{kr_c}} J_2(x_n) \quad (20)$$

Thus the final solution for the massive modes turns out to be

$$\psi^n(z_n) = \sqrt{kr_c} \frac{e^{2\sigma}}{e^{kr_c\pi}} \frac{J_2(z_n)}{J_2(x_n)} \quad (21)$$

We shall compare the lowest-lying massive modes of the dilaton field with those of graviton as well as the bulk KR field in the next section, when we determine the masses for the latter.

The solution for the massless dilaton mode can be obtained from equation (13) for  $m_n^d = 0$  and is given by,

$$\psi^0 = \frac{C_1}{4kr_c} e^{4\sigma} + C_2 \quad (22)$$

With the help of the same boundary conditions that has been used for the massive case, to determine the constants  $(C_1, C_2)$ , one finally arrives at the solution for the massless mode as,

$$\psi^0 = \sqrt{kr_c} \quad (23)$$

The complete solution for the dilaton field is then,

$$\phi(x, y) = \sqrt{k}\phi^0(x) + \sum_{n=1}^{\infty} \sqrt{k} \frac{e^{2\sigma}}{e^{kr_c\pi}} \frac{J_2(z_n)}{J_2(x_n)} \phi^n(x) \quad (24)$$

On our visible brane the field would appear as

$$\phi(x, y = \pi) = \sqrt{k}\phi^0(x) + \sum_{n=1}^{\infty} \sqrt{k} e^{kr_c\pi} \phi^n(x) \quad (25)$$

At this point it is useful to mention that the solution for the bulk massless dilaton is the same as that for the bulk graviton field. Unlike the massless case, the massive modes in both the cases appear with an exponential enhancement factor.

## 4 Bulk KR field coupled to dilaton

In this section we proceed with the term in the first square bracket in equation(11) where the dilaton field is coupled to the bulk KR field. Since we have already found out how the bulk dilaton field would appear on the visible brane (eqn(12)) and what would be the structure of the dilaton field (eqn(24)), we now use this solution in our subsequent analysis.

$$S_{KR} = \int d^4x \left[ \sum_n \sum_m \int dy e^{2\sigma} \chi^n \chi^m \left\{ -2e^{\frac{\sqrt{k}}{M^{3/2}}[\phi^0 + \sum_i \frac{e^{2\sigma} J_2(z_i)}{e^{kr_c\pi} J_2(x_i)} \phi^i(x)]} (\eta^{\mu\alpha} \eta^{\nu\beta} \eta^{\lambda\gamma} \partial_{[\mu} B_{\nu\lambda]}^n \partial_{[\alpha} B_{\beta\gamma]}^m \right. \right. \\ \left. \left. - 3 \left[ \frac{1}{r_c^2} \frac{e^{-2\sigma}}{\chi^m} \partial_y^2 \chi^m \right] \eta^{\mu\alpha} \eta^{\nu\beta} B_{\mu\nu}^n B_{\alpha\beta}^m \right\} \right] \quad (26)$$

$$= \int d^4x \left[ \sum_n \sum_m \int dy e^{2\sigma} \chi^n \chi^m \left\{ -2e^{\frac{\sqrt{k}}{M^{3/2}}[\phi^0 + \sum_i b_i(y) \phi^i(x)]} (H_{\mu\nu\lambda}^n H^{m\mu\nu\lambda} + 3m_m^2 B_{\mu\nu}^n B^{m\mu\nu}) \right\} \right] \quad (27)$$

As discussed earlier, the third rank anti-symmetric field strength  $H_{\mu\nu\lambda}^n = \partial_{[\mu} B_{\nu\lambda]}^n$  may be identified as torsion on our 4D brane,

$$b_i(y) = \frac{e^{2\sigma} J_2(z_i)}{e^{kr_c\pi} J_2(x_i)}$$

and

$$-\frac{1}{r_c^2} \frac{e^{-2\sigma}}{\chi^m} \partial_y^2 \chi^m = m_m^2$$

defines the mass of the  $m$ th mode of the KR field. This equation is much like equation (13).

In terms of  $z'_m = \frac{m_m^2}{k} e^\sigma$  the equation can be recast in the form

$$z'_m{}^2 \frac{d^2 \chi^m}{dz_m'^2} + z'_m \frac{d \chi^m}{dz_m'} + z'_m{}^2 \chi^m = 0 \quad (28)$$

This is a Bessel equation and solutions of this equation are Bessel functions of order 0.

$$\chi^m = \frac{1}{N'_m} [J_0(z'_m) + \alpha'_m Y_0(z'_m)] \quad (29)$$

where  $J_0(z'_m)$  and  $Y_0(z'_m)$  are Bessel and Neumann functions of order 0 and  $\alpha'_m$  and  $N'_m$  are two constants in this case to be determined from the boundary conditions. The continuity condition for the derivative of  $\chi^m$  at  $y = 0$  yields

$$\alpha'_m = -\frac{J'_0(\frac{m_m^t}{k})}{Y'_0(\frac{m_m^t}{k})} = -\frac{J_1(\frac{m_m^t}{k})}{Y_1(\frac{m_m^t}{k})} \quad (30)$$

Here  $J'_0(\frac{m_m^t}{k}) = \frac{d}{dz'_m} J_0(z'_m)|_{y=0}$ . Like the previous section we use the fact that  $e^{kr_c\pi} \gg 1$  and the mass values  $m_m^t$  on the brane are expected to be in the TeV range ( $\ll k$ ),

$$\alpha'_m = -\frac{\pi}{4} \left(\frac{m_m^t}{k}\right)^2 \ll 1 \quad (31)$$

The continuity condition for the derivative of  $\chi^m$  at  $y = \pi$  determines  $m_m^t$  from the roots of the equation

$$J_1(x'_m) = 0 \quad (32)$$

where  $x'_m = z'_m(\pi) = \frac{m_m^t}{k} e^{kr_c\pi}$ . It is interesting to see that the masses of the excited modes for both the KR field and dilaton field are governed by the same equation. While the masses of the higher modes of the dilaton field are the same as those of the graviton modes, the massive KR fields are scaled by a factor of  $\sqrt{3}$  from the other two. We enlist the masses of a few low-lying modes of graviton, dilaton and KR field in table 1.

$n$	1	2	3	4
$m_n^{grav} (TeV)$	1.66	3.04	4.40	5.77
$m_n^{dil} (TeV)$	1.66	3.04	4.40	5.77
$m_n^{KR} (TeV)$	2.87	5.26	7.62	9.99

Table 1: The masses of a few low-lying modes of the graviton, dilaton and Kalb-Ramond fields, for  $kr_c = 12$  and  $k = 10^{19} GeV$ .

So, as a result of these boundary conditions  $\chi^m$  looks like

$$\chi^m = \frac{1}{N'_m} J_0(z'_m) \quad (33)$$

where  $N'_m$  is yet to be determined. Now we go back to equation(27) and make a series expansion of the exponential term

$$\begin{aligned} S_{KR} = & \int d^4x \sum_n \sum_m \left[ \int dy e^{2\sigma} \chi^n \chi^m \left[ 1 + \frac{\sqrt{k}}{M^{3/2}} \left( \phi^0 + \sum_i b_i(y) \phi^i \right) \right. \right. \\ & \left. \left. + \frac{1}{2} \left( \frac{\sqrt{k}}{M^{3/2}} \right)^2 \left[ \phi^0 + \left( \sum_i b_i(y) \phi^i \right) \right]^2 + \dots \dots \right] \{ -2(H_{\mu\nu\lambda}^n H^{m\mu\nu\lambda} + 3m_m^t{}^2 B_{\mu\nu}^n B^{m\mu\nu}) \} \right] \quad (34) \end{aligned}$$



The first term of the series in the integral represent solely the lagrangian for the free KR field in 4D. The integral on the extra dimension, associated with this term, defines the orthogonality relation for the function  $\chi^n$  as,

$$\int dy e^{2\sigma} \chi^n \chi^m = \delta_{mn} \quad (35)$$

This is the same as found in ref[13]. The orthogonality relation determines the normalisation constant  $N'_m$  as

$$N'_n = \frac{e^{kr_c\pi}}{\sqrt{kr_c}} J_0(x'_n) \quad (36)$$

Using this the final solutions for the massive modes for the KR field turns out to be

$$\chi^n(z'_n) = \sqrt{kr_c} e^{-kr_c\pi} \frac{J_0(z'_n)}{J_0(x'_n)} \quad (37)$$

The solution for the massless mode of the KR field can be obtained by solving the differential equation for  $m_m^t = 0$ . This yields,

$$\chi^0 = C_3|y| + C_4 \quad (38)$$

The condition of self-adjointness implies the solution will be a constant which can be determined from the normalisation condition as,

$$\chi^0 = \sqrt{kr_c} e^{-kr_c\pi} \quad (39)$$

This identically matches our solution in ref [13] and thereby consolidates our claim that due to the presence of this large exponential suppression factor the KR field is heavily suppressed on the visible brane and thus makes the torsion imperceptible.

Let us now look back at equation (34). Different terms of the series exhibit different couplings (trilinear and higher order) between the KR field and dilaton field. The second term demonstrate a trilinear coupling of the form  $H^2\phi$ . The coupling of the massless mode of the dilaton field with the massless KR mode is suppressed by  $\frac{1}{M_P}$  (k is taken to be on the order of Planck scale). However, the trilinear coupling of the form  $\phi^0 H^0 H^n$  vanishes due to the orthogonality relation. We have to figure out the coupling scale of the massive modes of the dilaton field to the KR field (both for its massive and massless modes). But the subsequent terms of the series, representing quartic and higher order coupling between the two fields, are even more heavily suppressed by the factors of  $\frac{1}{M_P^2}$  or more. Thus there is hardly any phenomenological implication of such terms in the context of terrestrial experiments.

Next, we calculate the coupling of the massive dilaton fields with the KR fields on the brane. For this, we use the solutions which we have obtained so far for both the fields.

From equation (34), the strength of the trilinear coupling between the  $i$ th massive mode of the dilaton field and the  $n$ th and  $m$ th mode of the KR field is given by,

$$\frac{\sqrt{k}}{M^{3/2}} \int_{-\pi}^{\pi} dy e^{2\sigma} \chi^n \chi^m b_i(y) \quad (40)$$

Using the expression for  $b_i(y)$  and the solution for massive mode of KR field  $\chi^n$  as given in equation(37), the above integral becomes,

$$\frac{kr_c}{M_P e^{3kr_c\pi}} \int_{-\pi}^{\pi} dy e^{4\sigma} \frac{J_0(z'_n) J_0(z'_m) J_2(z_i)}{J_0(x'_n) J_0(x'_m) J_2(x_i)} \quad (41)$$

while the coupling expression for the massive dilaton field with the massless KR field is

$$\frac{kr_c}{M_P e^{3kr_c\pi}} \int_{-\pi}^{\pi} dy e^{4\sigma} \frac{J_2(z_i)}{J_2(x_i)} \quad (42)$$

We compute the integrals numerically taking  $kr_c = 12$  and considering the mass of different modes of the fields from table 1. The couplings for the all the massive fields turn out to be  $0.99 \times 10^{-3} \text{ GeV}^{-1}$  while the coupling for the massive dilaton with massless KR is  $7.51 \times 10^{-3} \text{ GeV}^{-1}$ . Thus the interaction strength is suppressed by  $\text{TeV}^{-1}$  and the lowest lying mass modes are also in the TeV scale. This yields the rather interesting possibility: the KR zero mode, which is shown to have extreme suppressed coupling with all matter fields, can now be produced in TeV-scale hadron collider experiments through the Drell-Yan process mediated by massive dilaton modes. Since the interaction with gluons or quark-antiquark pairs are now suppressed by only the TeV-scale, such production may have significant rates for sufficient integrated luminosity, and may show up in final states with missing transverse energy with some accompanying tagged jets. In addition, massive KR modes, too, can in principle be pair-produced, leading to final states that may look akin to those from graviton pairs, but may differ in angular distributions etc. due to the characteristic Lorenz structure of the interactions involved. Thus dilaton mediation proves to be of interest, so far as the experinmental signatures of the KR tower are concerned.

## 5 Conclusions

Extending our earlier work [13], we have considered the full bosonic sector of the low energy effective string action comprising of bulk dilaton, KR and gravity embedded in a five dimensional RS scenario, and have shown that the heavily suppressed nature of torsion on our visible brane is unaffected by its non-trivial coupling with the dilaton inside the bulk. This explanation of the torsion-free nature of our space-time originates in the warped braneworld model proposed by Randall and Sundrum.

We have also shown that such an warped extra dimensional model indicates the presence of new interactions between the Kaluza Klein modes of the dilaton and KR field which may be of significant phenomenological importance. The compactification of the fifth dimension gives rise to a spectrum of Kaluza Klein modes for all the fields including gravity. It is interesting to notice that the condition that determines the masses of the different Kaluza Klein modes of all the three fields are same, which eventually identifies the same mass for different modes for graviton and dilaton while the massive KR spectrum is scaled by a factor of  $\sqrt{3}$  with respect to the others[see table 1]. It is shown that the trilinear coupling between dilaton and KR field are of the order of  $\text{TeV}^{-1}$  and hence can lead to detectable signals in TeV scale accelerator experiments. The other higher order (quartic and so on) couplings are again very highly suppressed ( $\mathcal{O}(\frac{1}{M_P^2})$ ) and won't have a visible impact.

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